Cosmic Gamma Rays From Extragalactic Proton-Antiproton Annihilations

11-1 INTRODUCTION

The existence or nonexistence of antimatter in the universe is a question of importance in the fields of cosmology and particle physics (Alfvén, 1965). Recently, Harrison (1967) has proposed the existence of large amounts of antimatter in the universe in order to account for the condensation of matter into galaxies. The only practical way of determining the existence of such antimatter is by the detection of the γ-radiation produced when antimatter and matter annihilate into mesons having γ-ray-producing decay modes. The γ-ray spectra from such annihilations are presented here. They are based on the results of recent accelerator experiments on antiproton annihilation as discussed in chapter 3 and also include cosmological distortions of local annihilation γ-ray spectra.

11-2 THE LOCAL ANNIHILATION GAMMA-RAY SPECTRUM

The annihilation γ-ray spectrum (AGS) for $p-\bar{p}$ interactions involving kinetic energies less than 286 MeV (the threshold for nonannihilation pion production) is given by the expression

$$I_A(E_\gamma) = \frac{B}{4\pi} \int dv\, v f(v) \sum_s \int dE_s\, \sigma_s(E_s|v) \sum_d \zeta_{\gamma d} R_{\gamma d} f_{ds}(E_\gamma|E_s) \quad (11\text{–}1)$$

as shown in chapter 3, where v is the relative velocity of the proton and antiproton, $f(v)$ is the normalized distribution function over this velocity, s is an index representing the particular type of particle produced in the annihilation, $\sigma_s(E_s|v)$ is the cross section times multiplicity of particles of type s and energy E_s produced in a collision of velocity v, and the index d specifies a specific decay mode with a branching ratio $R_{\gamma d}$ which produces $\zeta_{\gamma d}$ γ-rays with a normalized energy distribution $f(E_\gamma|E_s)$.

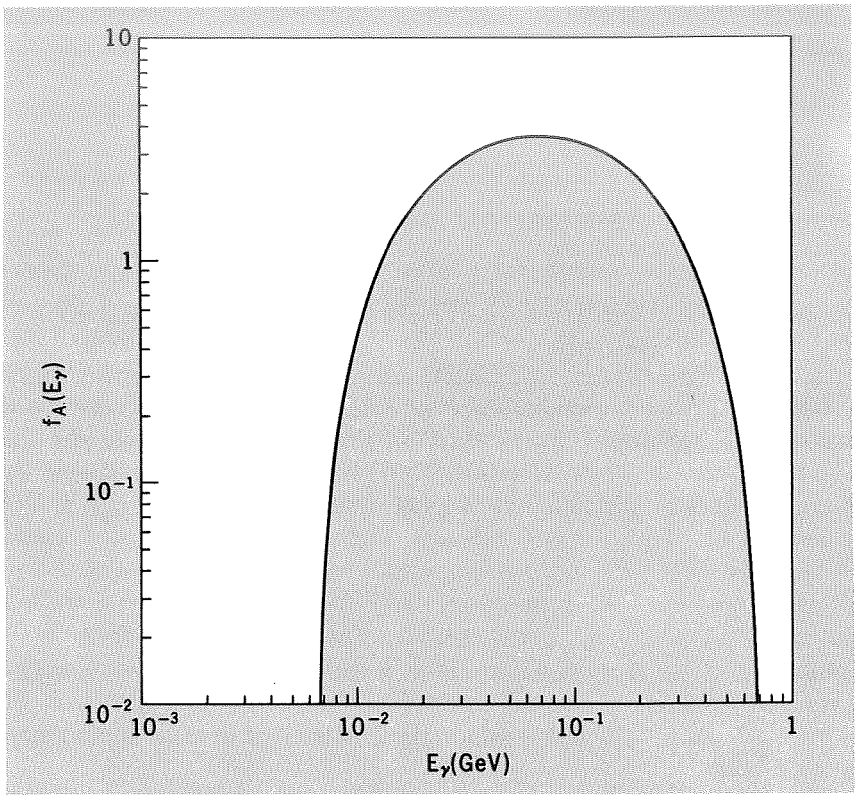

FIGURE 11–1.—The normalized local differential γ-ray spectrum from $p-\bar{p}$ annihilation.

The quantity B is defined as the product of interacting proton and antiproton densities integrated over a line-of-sight pathlength l; i.e.,

$$B = \int_0^{l_{\max}} dl\; n_p(l) n_{\bar{p}}(l) \tag{11–2}$$

Based on the calculations of Morgan and Hughes (1969), we showed in chapter 3 that $(E_s|v)$ can be represented as having a power-law velocity dependence using three different exponents given by equation (3–31) for three different temperature ranges given in equation (3–30). Using equations (3–47) and (3–48), we may thus write

$$I_A(E_\gamma) = \frac{B}{4\pi}\, \sigma_{\text{III}} c \zeta_\gamma \left[\frac{2}{\pi^{1/2}}\, \Gamma(1.68) \left(\frac{2kT}{mc^2}\right)^{0.18}\right] f_A(E_\gamma)$$
$$\text{for } 10\text{ K} \leqslant T \leqslant 10^4\text{ K} \tag{11–3}$$

and

$$I_A(E_\gamma) = \frac{B}{4\pi}\,\sigma_\Pi c\zeta_\gamma \left(\frac{2mc^2}{\pi kT}\right)^{1/2} f_A(E_\gamma) \qquad \text{for } 10^4\ \text{K} \lesssim T \lesssim 10^{11}\ \text{K} \quad (11\text{–}4)$$

The local annihilation spectrum $f_A(E_\gamma)$, as calculated in chapter 3, is shown in figure 11–1.

11-3 REDSHIFTED SPECTRA

For cosmological γ-ray production

$$B \rightarrow \int dz\, n_p(z) n_{\bar{p}}(z)\, \frac{dl}{dz} \quad (11\text{–}5)$$

where, from equation (9–41),

$$\frac{dl}{dz} = \frac{c}{H_0(1+z)^2(1+\Omega z)^{1/2}}$$

Taking (from equation 9–38)

$$n_p(z) = n_{p,0}(1+z)^3$$

and (11–6)

$$n_{\bar{p}}(z) = n_{\bar{p},0}(1+z)^3$$

and defining

$$B_0 \equiv \left(\frac{c}{H_0}\right) n_{p,0} n_{\bar{p},0} \quad (11\text{–}7)$$

we find by using equation (9–21) that

$$I_A(E_\gamma) = \frac{B_0 \sigma_i c^{\delta_i}}{2\pi^{3/2}} \left(\frac{2k}{m}\right)^{\left(\frac{1-\delta_i}{2}\right)} \Gamma\left(2 - \frac{\delta_i}{2}\right)$$

$$\times \int dz [T(z)]^{\left(\frac{1-\delta_i}{2}\right)} \frac{(1+z)^2}{(1+\Omega z)^{1/2}}\, G_A[(1+z)E_\gamma] e^{-\tau(E_\gamma, z)} \quad (11\text{–}8)$$

where $i = \text{II, III}$; $\sigma_{\text{II}} = 2$; $\delta_{\text{III}} = 0.64$; $\sigma_{\text{II}} = 1.1 \times 10^{-27}\ \text{cm}^2$; $\sigma_{\text{III}} = 2.6 \times 10^{-18}\ \text{cm}^2$; $G_A(E_\gamma) \equiv \zeta_\gamma f_A(E_\gamma)$ where $f_A(E_\gamma)$ is shown in figure 11–1; and T_m is the matter temperature of the universe, which is a function of z.

At the earliest stages in the evolution of the universe, when the matter in the universe was in thermal equilibrium with the universal blackbody radiation due to Compton interactions between thermal photons and electrons, T_m was equal to T_r (the temperature of the radiation field). Zeldovich, Kurt, and Syunyaev (1969) have shown that even though the intergalactic gas has cooled to the 50-percent neutral point by the time corresponding to a redshift of 1200 to 1300, the small fraction of ionized material left at lower redshifts is enough to sustain temperature equilibrium between matter and the radiation field until a redshift of 150 to 200.

The photon density of a radiation field at temperature T_0 is given from the Planck formula as

$$n_{r,0}(\epsilon, T_0)\,d\epsilon = \frac{8\pi}{h^3c^3}\,\frac{\epsilon^2\,d\epsilon}{\exp\left[\dfrac{\epsilon}{kT_{r,0}}\right]-1} \tag{11-9}$$

At a redshift $z \neq 0$, this distribution is given by

$$n_r(z, \epsilon)\,d\epsilon = \frac{8\pi}{h^3c^3}\,\frac{(1+z)^3\epsilon^2\,d\epsilon}{\exp\left[\dfrac{(1+z)\epsilon}{kT_{r,0}}\right]-1}$$

$$= (1+z)^3 n_{r,0}[\epsilon, T_r(z)]\,d\epsilon \tag{11-10}$$

where

$$T_r(z) \equiv (1+z)T_{r,0} \tag{11-11}$$

Thus, we find that the redshifted Planckian maintains its form with the parameter $T(z)$ in the exponent given by equation (11–11). The photon density redshifts in the same manner as matter density as given by equation (10–1). The energy density in the radiation field then redshifts as

$$\rho_r = \rho_{r,0}(1+z)^4 \tag{11-12}$$

in accord with the definition of $T_r(z)$ given by equation (11–11) and the relation $\rho \sim T^4$.

We therefore find that

$$T_m(z) = T_r(z) = T_{r,0}(1+z) \qquad \text{for } z \gtrsim 150 \text{ to } 200 \tag{11-13}$$

At lower redshifts, when the matter has thermally decoupled from the radiation field, the momentum distribution of the atoms in the intergalactic gas is given by the Maxwellian,

$$n_0(p, T_{m,0})\,dp = \text{const} \times p^2 e^{-p^2/2mkT_{m,0}}\,dp \tag{11-14}$$

These atoms lose momentum through collisions due to the effect of the overall adiabatic expansion of the universe; their resultant momentum change is

$$p(z) = p_0 \frac{R_0}{R(z)} = p_0(1+z) \tag{11-15}$$

Substituting equation (11–15) into equation (11–14), we then find

$$n_m(z, p)dp = (1+z)^3 n_{m,0}[p, T_m(z)]dp \tag{11-16}$$

with

$$T_m(z) = T_{m,0}(1+z)^2 \tag{11-17}$$

From equation (11–16), it is immediately evident that the total number of atoms, $\sim nR^3$, is conserved during the expansion, as must be the case.

We thus find that

$$T_m(z) = T_{m,0}(1+z)^2 \qquad \text{for } z \lesssim 150 \text{ to } 200 \tag{11-18}$$

with

$$T_m(150 \text{ to } 200) \simeq T_r(150 \text{ to } 200) \tag{11-19}$$

The adiabatic expansion concept may be used to link the derivations of equations (11–11) and (11–17) for the temperature-redshift relations. In an adiabatic expansion, it is well known that

$$TV^{(\gamma-1)} = \text{constant} \tag{11-20}$$

where the quantity γ here is the ratio of specific heat at constant pressure to that at constant temperature and V is the volume of the gas (in this case the volume of the universe) which from equation (9–15) is given by

$$V(z) = V_0(1+z)^{-3} \tag{11-21}$$

For photons and monatomic gases at relativistic temperatures ($pc \simeq E$), $\gamma = 4/3$, whereas for nonrelativistic gases, $\gamma = 5/3$. Thus, from equations (11–21) and (11–20) we immediately obtain equation (11–11) for photons and relativistic gases and equation (11–17) for nonrelativistic gases.

Equations (11–11) and (11–17) may then be used in equation (11–8) to evaluate $I_A(E_\gamma)$.

It is of interest to note that since $f_A(E_\gamma)$ is nonvanishing only within a restricted energy range (given by eq. (3–76)), there exists an energy region

$$E_\gamma \lesssim 5 \text{ MeV} \tag{11-22}$$

TABLE 11–1.—*Exponents of Cosmological Power-Law Annihilation γ-Ray Spectra*

Ω	z	$E_{\gamma 0} \simeq \frac{70 \text{ MeV}}{(1+z)}$	δ	$\alpha+1$
1	0–150	500 keV–70 MeV	0.64	2.86
	150–1000	70 keV–500 keV	0.64	2.68
	10^3–10^8	< 70 keV	2	2.00
$\simeq 0$	0–150	500 keV–70 MeV	0.64	3.36
	150–1000	70 keV–500 keV	0.64	3.18
	10^3–10^8	< 70 keV	2	2.50

where the integral in equation (11–8) is independent of the form of $f_A(E_\gamma)$ and the resultant spectrum has a power-law form when absorption may be neglected. In actuality, the power-law form of the spectrum is a valid approximation even over a wider energy range, where a change of the limits of integration over $f_A(E_\gamma)$ does not significantly affect the value of the integral. In this case, we find

$$I_A(E_\gamma) = 2B_0\sigma_\pi\zeta_\pi cKE_\gamma^{-(\alpha+1)} \sim E_\gamma^{-(\alpha+1)} \tag{11–23}$$

where

$$K \equiv \int_{5 \text{ MeV}}^{919 \text{ MeV}} dE_\gamma E_\gamma^\alpha f_\gamma(E_\gamma) \tag{11–24}$$

and α is the exponent of the total power-law dependence of the integrand on $(1+z)$; i.e.,

$$I_A(E_\gamma) \sim \int dz(1+z)^\alpha f_A[(1+z)E_\gamma] \tag{11–25}$$

For an Einstein-de Sitter universe, $\Omega = 1$, whereas for a low-density Friedman universe, $\Omega \simeq 0$. Thus, we find the power-law production spectra with exponents given in table 11–1 for the various ranges of γ-ray energies at $z = 0$.